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Constraining the *p*-Mode–*g*-Mode Tidal Instability with GW170817

B. P. Abbott *et al.*[*]
(LIGO Scientific Collaboration and Virgo Collaboration)

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We analyze the impact of a proposed tidal instability coupling p modes and g modes within neutron stars on GW170817. This nonresonant instability transfers energy from the orbit of the binary to internal modes of the stars, accelerating the gravitational-wave driven inspiral. We model the impact of this instability on the phasing of the gravitational wave signal using three parameters per star: an overall amplitude, a saturation frequency, and a spectral index. Incorporating these additional parameters, we compute the Bayes factor ($\ln B^{pg}_{!pg}$) comparing our p-g model to a standard one. We find that the observed signal is consistent with waveform models that neglect p-g effects, with $\ln B^{pg}_{!pg} = 0.03^{+0.70}_{-0.58}$ (maximum *a posteriori* and 90% credible region). By injecting simulated signals that do not include p-g effects and recovering them with the p-g model, we show that there is a $\simeq$50% probability of obtaining similar $\ln B^{pg}_{!pg}$ even when p-g effects are absent. We find that the p-g amplitude for 1.4 $M_\odot$ neutron stars is constrained to less than a few tenths of the theoretical maximum, with maxima *a posteriori* near one-tenth this maximum and p-g saturation frequency ~70 Hz. This suggests that there are less than a few hundred excited modes, assuming they all saturate by wave breaking. For comparison, theoretical upper bounds suggest $\lesssim 10^3$ modes saturate by wave breaking. Thus, the measured constraints only rule out extreme values of the p-g parameters. They also imply that the instability dissipates $\lesssim 10^{51}$ erg over the entire inspiral, i.e., less than a few percent of the energy radiated as gravitational waves.

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Introduction.—Detailed analysis of the gravitational-wave (GW) signal received from the first binary neutron star (NS) coalescence event (GW170817 [1]) constrains the tidal deformability of NSs and thus the equation of state (EOS) above nuclear saturation density [2–4]. Studies of NS tidal deformation typically focus on the linear, quasi-static tidal bulge induced in each NS by its companion. Such deformations modify the system's binding energy and GW luminosity and thereby alter its orbital dynamics. The degree of deformation is often expressed in terms of the tidal deformability $\Lambda_i \propto (R_i/m_i)^5$ of each component [5], or a particular mass-weighted average thereof ($\tilde{\Lambda}$) [2]. The strong dependence on compactness R/m means that a stiffer EOS, which has larger R for the same m, imprints larger tidal signals than a softer EOS. Current analyses of GW data from the LIGO [6] and Virgo [7] detectors favor a soft EOS [3,8]. Specifically, Ref. [2] finds $\tilde{\Lambda} \lesssim 730$ at the 90% credible level for all waveform models considered, allowing for the components to spin rapidly. The pressure at twice nuclear saturation density is also constrained to $P = 3.5^{+2.7}_{-1.7} \times 10^{34}$ dyn/cm^2 (median and 90% credible region) [3] assuming small component spins. In addition to GW phasing, the EOS dependence of $\tilde{\Lambda}$ should correlate with postmerger signals [9], possible tidal disruptions, and kilonova observations [10]. Observed light curves for the kilonova suggest a lower bound of $\tilde{\Lambda} \gtrsim 200$ [11,12].

Although some dynamical tidal effects are incorporated in these analyses (see, e.g., Refs. [2,13]), the impact of several types of dynamical tidal effects are neglected because they are believed to be small or have large theoretical uncertainties. These effects arise because tidal fields, in addition to raising a quasistatic bulge, excite stellar normal modes. Three such excitation mechanisms are (i) resonant linear excitation, (ii) resonant nonlinear excitation, and (iii) nonresonant nonlinear excitation (see, e.g., Ref. [14]). The first occurs when the GW frequency (the oscillation frequency of the tidal field) sweeps through a mode's natural frequency (see, e.g., Refs. [15–22]). However, since the GW frequency increases rapidly during the late inspiral, the time spent near resonance is too short to excite modes to large amplitudes. As a result, for modes with natural frequencies within the sensitive bands of ground-based GW detectors, the change in orbital phasing is expected to be small ($\Delta\Psi \lesssim 10^{-2}$ rad) unless the stars are rapidly rotating [17–19]. The impact of resonant nonlinear mode excitation (i.e., the parametric subharmonic instability) is likewise limited by the swiftness of the inspiral [23].

The proposed p-g tidal instability is a nonresonant, nonlinear instability in which the tidal bulge excites a low-frequency buoyancy-supported g mode and a high-frequency pressure-supported p mode [23–26]. It occurs in the inner core of the NS, where the stratification is weak and the shear due to the tidal bulge is especially susceptible to instability. Unlike resonantly excited modes, an unstable

[*]Full author list given at the end of the Letter.

p-g pair continuously drains energy from the orbit once excited, even after the orbital frequency changes significantly. There are many potentially unstable p-g pairs, each becoming unstable at a different frequency and growing at a different rate. Although there is considerable uncertainty about the number of unstable pairs, their exact growth rates, and how they saturate, estimates suggest that the impact could be measurable with current detectors [27].

In this Letter, we investigate the possible impact of the p-g instability on GW170817 using the phenomenological model developed in Ref. [27]. The model describes the energy dissipated by the instability within each NS, indexed by i, in terms of three parameters: (i) an overall amplitude A_i, which is related to the number of modes participating in the instability, their growth rates, and their saturation energies, (ii) a frequency f_i corresponding to when the instability saturates, and (iii) a spectral index n_i describing how the saturation energy evolves with frequency. In the section "Phenomenological model," we describe our models in detail. In the section "Model selection," we compare the statistical evidence for models that include the p-g instability relative to those that do not. In the section "Parameter inference," we investigate the constraints on the p-g parameters from GW170817, and in the section "Discussion," we conclude.

Phenomenological model.—Following Ref. [27], we extend a post-Newtonian (PN) waveform by including a parametrized model of the p-g instability. For the initial PN model, we use the TaylorF2 frequency-domain approximant (see, e.g., Ref. [28]) terminated at the innermost stable circular orbit, which includes the effects of linear tides ($\tilde{\Lambda}$) and spins aligned with the orbital angular momentum (the impact of misaligned spins on p-g effects is not known). Waveform systematics between different existing approximants may be important for small p-g effects. However, by comparing the waveform mismatches between several other models (TaylorF2, SEOBNRT, PhenomDNRT, and PhenomPNRT, see Ref. [2]), we find these systematics induce waveform mismatches that correspond to p-g phenomenological amplitudes roughly an order of magnitude smaller than the upper limits set by our analysis (see section "Parameter inference"). We expect the TaylorF2 approximant to be reasonably accurate and defer a complete analysis of waveform systematics to future work.

Assuming the p-g effects are a perturbation to the TaylorF2 approximant, we find that they modify the phase in the frequency domain by

$$\Delta\Psi(f) = -\frac{2C_1}{3B^2(3-n_1)(4-n_1)}\left\{\Theta_1\left(\frac{f}{f_{\mathrm{ref}}}\right)^{n_1-3} + (1-\Theta_1)\left(\frac{f_1}{f_{\mathrm{ref}}}\right)^{n_1-3}\left[(4-n_1)-(3-n_1)\left(\frac{f}{f_1}\right)\right]\right\} + (1 \leftrightarrow 2), \quad (1)$$

where f_i is the saturation frequency, $f_{\mathrm{ref}} \equiv 100$ Hz is a reference frequency with no intrinsic significance, $C_i = [2m_i/(m_1+m_2)]^{2/3}A_i$, $B = (32/5)(G\mathcal{M}\pi f_{\mathrm{ref}}/c^3)^{5/3}$, $\mathcal{M} = (m_1 m_2)^{3/5}/(m_1+m_2)^{1/5}$, and $\Theta_i = \Theta(f - f_i)$ where Θ is the Heaviside function. This approximant is slightly different than that of Ref. [27] because they incorrectly applied the saddle-point approximation to obtain the frequency-domain waveform from time-domain phasing [29]. This correction renders the p-g instability slightly more difficult to measure than predicted in Ref. [27], although the observed behavior is qualitatively similar. Specifically, we find that in order to achieve the same $|\Delta\Psi|$, A_i needs to be larger than Ref. [27] found by a factor of $\sim(4-n_i)$, although the precise factor also depends on the other p-g parameters.

The $\Delta\Psi$ expression contains three types of terms: a constant term, a linear term $\propto (1-\Theta_i)f$, and a power-law term $\propto \Theta_i f^{n_i-3}$. The constant term corresponds to an overall phase offset and is degenerate with the orbital phase at coalescence. The linear term corresponds to a change in the time of coalescence; because the p-g instability transfers energy from the orbit to stellar normal modes, the binary inspirals faster than it would if the effect was absent. The power-law term accounts for the competition between the rate of p-g energy dissipation and the rate of inspiral, both of which increase as f increases. As argued in Ref. [27], we expect $n_i < 3$, which implies that the phase shift accumulates primarily at frequencies just above the "turn-on" (saturation) frequency $f \gtrsim f_i$.

When $n_i < 3$, p-g effects are most important at lower frequencies whereas linear tides ($\tilde{\Lambda}$) and spins ($\chi_i = cS_i/Gm_i^2$, where S_i is the spin-angular momentum of each component) have their largest impact at higher frequencies (see, e.g., Ref. [30]). The priors placed on the latter quantities can, however, affect our inference of p-g parameters.

In order to account for a possible dependence on the component masses (m_i), we parametrize our model using a Taylor expansion in the p-g parameters around $m_i = 1.4\ M_\odot$ and sample from the posterior using the first two coefficients. Our model computes A_i as

$$A_i(m_i) = A_0 + \left(\frac{dA}{dm}\bigg|_{1.4\ M_\odot}\right)(m_i - 1.4\ M_\odot), \quad (2)$$

and uses A_0 and dA/dm instead of A_1 and A_2. The model uses similar representations for f_i and n_i in terms of the parameters f_0, df/dm, n_0, and dn/dm. We assume a uniform prior on $\log_{10} A_0$ within $10^{-10} \le A_0 \le 10^{-5.5}$, a uniform prior in f_0 within 10 Hz $\le f_0 \le$ 100 Hz, and a uniform prior in n_0 within $-1 \le n_0 \le 3$. The priors on the first-order terms ($dA/dm, df/dm, dn/dm$) are the same as those in Ref. [27]; when $m_1 \sim m_2$, they imply $A_1 \sim A_2$, etc.

We investigate GW170817 using data from several different frequency bands and with different spin priors, but unless otherwise noted we focus on results for data above 30 Hz with $|\chi_i| \leq 0.89$. Throughout this Letter, results from GW170817 were obtained using the same data conditioning as Ref. [2], including the removal of a short-duration noise artifact from the Livingston data (Ref. [31] and discussion in Ref. [1]) along with other independently measured noise sources (see, e.g., Refs. [32–35]), calibration [36,37], marginalization over calibration uncertainties, and whitening procedures [38,39]. We use the publicly available LALInference software package throughout [40,41].

Model selection.—Using GW data from GW170817, we perform Bayesian model selection. We compare a model that includes linear tides, spin components aligned with the orbital angular momentum, and PN phasing effects up to 3.5 PN phase terms ($\mathcal{H}_{!pg}$) to an extension of this model that also includes p-g effects ($\mathcal{H}_{pg}$). Since we have nested models ($\mathcal{H}_{!pg}$ is obtained from $\mathcal{H}_{pg}$ as $A_i \to 0$), we use the Savage-Dickey density ratio (see, e.g., Refs. [42–44]) to estimate the Bayes factor ($B^{pg}_{!pg} = p(D|\mathcal{H}_{pg})/p(D|\mathcal{H}_{!pg})$, where D refers to the observed data). Because we use a uniform-in-$\log_{10} A_0$ prior, $\mathcal{H}_{pg}$ does not formally include $A_i = 0$. Nonetheless, our lower limit on A_i is sufficiently small that $\mathcal{H}_{!pg}$ is effectively nested in $\mathcal{H}_{pg}$. Specifically, we sample from the model's posterior distribution [40,41] and calculate

$$\begin{aligned}\lim_{A_i\to 0}\left[\frac{p(A_i|D,\mathcal{H}_{pg})}{p(A_i|\mathcal{H}_{pg})}\right] &= \lim_{A_i\to 0}\left[\frac{1}{p(D|\mathcal{H}_{pg})}\int d\theta df_i dn_i p(D|\theta,A_i,f_i,n_i;\mathcal{H}_{pg})p(\theta|\mathcal{H}_{pg})p(f_i,n_i|A_i,\mathcal{H}_{pg})\right]\\ &= \frac{1}{p(D|\mathcal{H}_{pg})}\int d\theta p(D|\theta;\mathcal{H}_{!pg})p(\theta|\mathcal{H}_{!pg})\left[\frac{p(\theta|\mathcal{H}_{pg})}{p(\theta|\mathcal{H}_{!pg})}\right]\int df_i dn_i p(f_i,n_i|A_i,\mathcal{H}_{pg})\\ &= \frac{p(D|\mathcal{H}_{!pg})}{p(D|\mathcal{H}_{pg})}\left\langle\frac{p(\theta|\mathcal{H}_{pg})}{p(\theta|\mathcal{H}_{!pg})}\right\rangle_{p(\theta|D,\mathcal{H}_{!pg})},\end{aligned} \tag{3}$$

where θ refers to all parameters besides the p-g phenomenological parameters; we note that $\int df dn p(f_i,n_i|A_i,\mathcal{H}_{pg}) = 1\,\forall A_i$, and $\langle x\rangle_p$ denotes the average of x with respect to the measure defined by p. Assuming that $p(\theta|\mathcal{H}_{pg}) = p(\theta|\mathcal{H}_{!pg})$, we determine $\ln B^{pg}_{!pg}$ from the ratio, as $A_i \to 0$, of the marginal distribution of A_i *a priori* to the distribution *a posteriori*:

$$\ln B^{pg}_{!pg} = \lim_{A_i\to 0}[\ln p(A_i|\mathcal{H}_{pg}) - \ln p(A_i|D,\mathcal{H}_{pg})]. \tag{4}$$

This allows us to directly measure $\ln B^{pg}_{!pg}$ by extracting $p(A|D,\mathcal{H}_{pg})$ from Monte Carlo analyses with a known prior $p(A|\mathcal{H}_{!pg})$. We confirmed that this estimate agrees with estimates from both nested sampling [45] and thermodynamic integration [46].

Figure 1 shows $\ln B^{pg}_{!pg}$ as a function of $f_{\rm low}$, the minimum GW frequency considered. At a given $f_{\rm low}$, we show the distribution of $\ln B^{pg}_{!pg}$ due to the sampling uncertainty from the finite length of our MCMC chains. The solid and dashed curves correspond to the high-spin ($|\chi_i| \leq 0.89$) and low-spin ($|\chi_i| \leq 0.05$) priors discussed in Refs. [1–3].

For certain combinations of $f_{\rm low}$ and $|\chi_i|$, we find $\ln B^{pg}_{!pg} > 0$, suggesting $\mathcal{H}_{pg}$ is more likely than $\mathcal{H}_{!pg}$. In order to assess how likely such values are, we calculate $\ln B^{pg}_{!pg}$ for a large number of simulated, high-spin signals with $A_i = 0$ and distinct realizations of detector noise from times near GW170817. We find that simulated signals without p-g effects can readily produce $\ln B^{pg}_{!pg}$ at least as large as the ones we measured from GW170817. For example, for the 30 Hz high-spin data we obtain $\ln B^{pg}_{!pg} = 0.03^{+0.70}_{-0.58}$ (maximum *a posteriori* and 90% credible region; bottom panel of Fig. 1), whereas approximately half of our simulated signals yield $\ln B^{pg}_{!pg}$ at least this large, i.e., a false alarm probability (FAP) ≈50%. We focus on the 30 Hz, high-spin data because

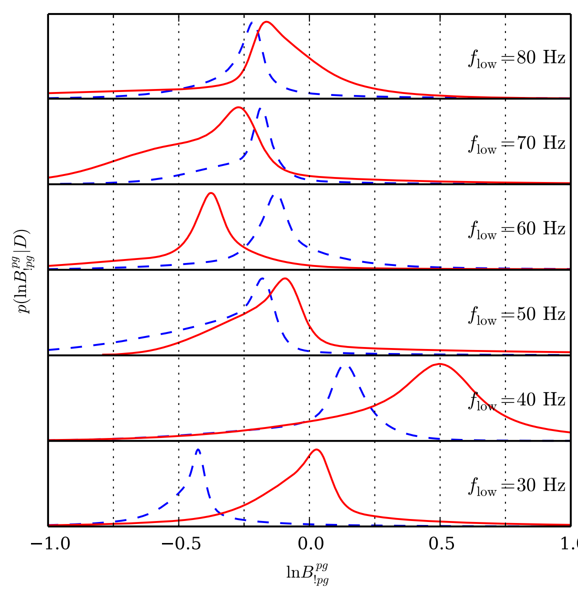

FIG. 1. Distributions of $\ln B^{pg}_{!pg}$ due to sampling uncertainty when analyzing GW170817 data with different values of $f_{\rm low}$. The solid red curves assume high-spin priors ($|\chi_i| \leq 0.89$) and the dashed blue curves assume low-spin priors ($|\chi_i| \leq 0.05$).

it corresponds to the largest bandwidth investigated and the largest signal-to-noise ratio. The high-spin prior is the most inclusive prior considered, and therefore allows the most model freedom when fitting p-g effects.

In our model of the instability, the phase shift $\Delta\Psi$ accumulates primarily at frequencies just above the saturation frequency $f \gtrsim f_i$. Therefore, if it is present, its impact should become more apparent as we decrease the minimum GW frequency considered from $f_{\rm low} \gg f_i$ to $f_{\rm low} \lesssim f_i$. We do see some indication of this behavior in Fig. 1. However, we note that if our phenomenological model breaks down at $f < f_i$ due to poor modeling of the presaturation behavior (e.g., if our step-function turn-on at f_i is not a good approximation to the instability's induced phase shift), we might expect $\ln B^{pg}_{!pg}$ to decrease as we lower $f_{\rm low}$ below f_i. If the fidelity of our model is sufficiently poor, we could be insensitive to p-g effects even at frequencies above $f_{\rm low}$.

Parameter inference.—We now investigate the constraints obtained from GW170817. Figure 2 shows the joint posterior distributions for both $\mathcal{H}_{pg}$ and $\mathcal{H}_{!pg}$. We find that $\mathcal{H}_{pg}$ and $\mathcal{H}_{!pg}$ yield similar posterior distributions for all non-p-g parameters, including both extrinsic and intrinsic parameters. The constraints on the chirp mass ($\mathcal{M}$), effective spin $\chi_{\rm eff} = (m_1\chi_1 + m_2\chi_2)/(m_1 + m_2)$, and $\tilde{\Lambda}$ are slightly weaker in $\mathcal{H}_{pg}$ than $\mathcal{H}_{!pg}$. This is because $\mathcal{H}_{pg}$ provides extra freedom to the signal's duration in the time domain.

Regarding the p-g parameters, we find a noticeable peak near $A_0 \sim 10^{-7}$ with a flat tail to small A_0. We find $A_0 \leq 3.3 \times 10^{-7}$ assuming a uniform-in-$\log_{10} A_0$ prior and $A_0 \leq 6.8 \times 10^{-7}$ assuming a uniform-in-A_0 prior, both at 90% confidence. The upper limit with a uniform-in-A_0 prior is larger only because we weight larger values of A_0 more *a priori* than with a uniform-in-$\log_{10} A_0$ prior. We also find a peak at $f_0 \sim 70$ Hz. The peaks persist when we analyze the data from each interferometer separately, with reasonably consistent locations and shapes (Fig. 2). However, we find that the simulated signals with $A_i = 0$ can produce similar peaks, suggesting they may be due to noise alone. Similar to Ref. [27], we find that n_i is not strongly constrained and the gradient terms in the Taylor expansions are not measurable.

Theoretical arguments suggest an upper bound of $A_0 \lesssim 10^{-6}$ [27]. Therefore, our A_0 constraint only rules out the most extreme values of the p-g parameters.

Discussion.—While GW170817 is consistent with models that neglect p-g effects, it is also consistent with a broad range of p-g parameters. The constraints from GW170817 imply that there are $\lesssim$200 excited modes at $f = 100$ Hz, assuming all modes grow as rapidly as possible and saturate at their breaking amplitudes ($\lambda = \beta = 1$ in Eq. (7) of Ref. [27]) and that the frequency at which modes become unstable is well approximated by f_0. For comparison, theoretical arguments suggest an upper bound of $\sim10^3$

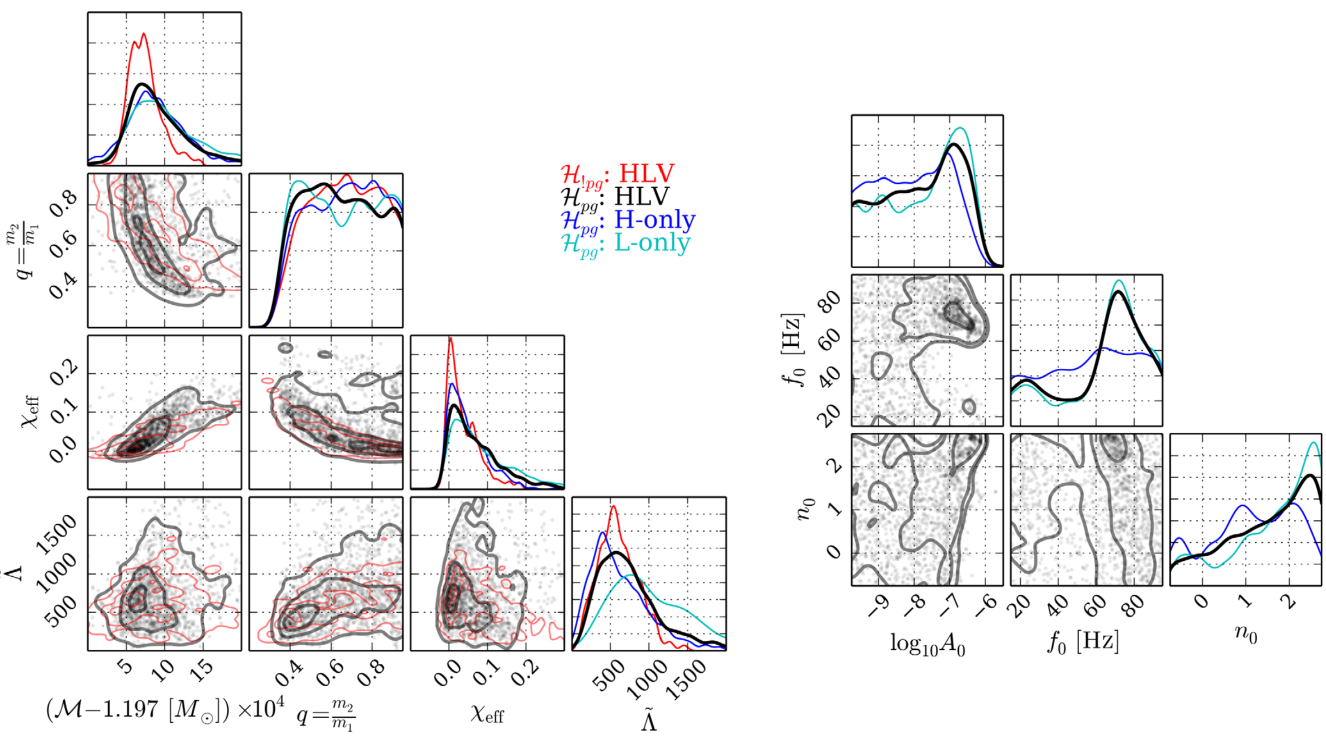

FIG. 2. Posterior distributions for $\mathcal{H}_{!pg}$ (red) and $\mathcal{H}_{pg}$ with Hanford, Livingston, and Virgo data (thick black, gray shading), Hanford data only (dark blue), and Livingston data only (light blue) using GW data above 30 Hz, $|\chi_i| \leq 0.89$, and a uniform-in-$\log_{10} A_0$ prior. Left: a subset of parameters shared by $\mathcal{H}_{!pg}$ and $\mathcal{H}_{pg}$. Right: a subset of parameters belonging only to $\mathcal{H}_{pg}$. We only show one-dimensional posteriors for the single instrument data, although the multidimensional posteriors are similarly consistent with the full $\mathcal{H}_{pg}$ data. Contours in the two-dimensional distributions represent 10%, 50%, and 90% confidence regions.

modes saturating by wave breaking [27]. More modes may be excited if they grow more slowly or saturate below their wave breaking energy.

We can also use the measured constraints to place upper limits on the amount of energy dissipated by the p-g instability. As Fig. 3 shows, p-g effects dissipate $\lesssim 2.7 \times 10^{51}$ erg throughout the entire inspiral at 90% confidence. In comparison, GWs carry away $\gtrsim 10^{53}$ erg. This implies time-domain phase shifts $|\Delta\phi| \lesssim 7.6$ rad (0.6 orbits) at 100 Hz and $|\Delta\phi| \lesssim 32$ rad (2.6 orbits) at 1000 Hz after accounting for the joint uncertainty in component masses, spins, linear tides, and p-g effects.

A g mode with natural frequency f_g is predicted to become unstable at a frequency $f_{\rm crit} \simeq 45\ {\rm Hz}(f_g/10^{-4}\lambda f_{\rm dyn})^{1/2}$, where $f_{\rm dyn}$ is the dynamical frequency of the NS and λ is a slowly varying function typically between 0.1–1 [25,27]. Since the modes grow quickly, the frequency at which the instability saturates is likely close to the frequency at which the modes become unstable ($f_0 \simeq f_{\rm crit}$). If we assume that the observed peak near $f_0 \sim 70$ Hz is not due to noise alone, then the maximum *a posteriori* estimate for f_0 along with approximate values for the masses ($1.4\ M_\odot$) and radii (11 km) of the components [3] imply $f_g \simeq 0.5$ Hz.

With several more signals comparable to GW170817, it should be possible to improve the amplitude constraint to $A_0 \lesssim 10^{-7}$. Obtaining even tighter constraints will likely require many more detections, especially since most events will have smaller SNR. Future measurements will also benefit from a better understanding of how the instability saturates. To date, there have only been detailed theoretical studies of the instability's threshold and growth rate [23–26], not its saturation. As a result, we cannot be certain of the fidelity of our phenomenological model.

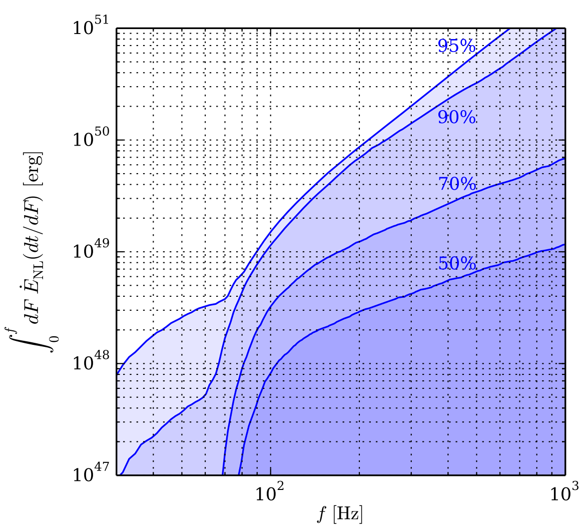

FIG. 3. Upper limits on the cumulative energy dissipated by the p-g instability as a function of frequency. We note the relatively strong constraints at lower frequencies where p-g effects are more pronounced.

While this Letter was in review, related work was posted [47] with the conclusion that the $\mathcal{H}_{!pg}$ model is strongly favored over the $\mathcal{H}_{pg}$ model by a factor of at least 10^4. In Ref. [48], some of the authors of this work investigate the origin of the discrepancy by analyzing publicly available posterior samples from Ref. [47]. Contrary to the claims in Ref. [47], they find that samples from Ref. [47] yield $B^{pg}_{!pg} \sim 1$ and therefore conclude that their posterior data, like what is presented here, do not disfavor the $\mathcal{H}_{pg}$ model. Reference [48] suggests that the error stems from using too few temperatures when implementing thermodynamic integration.

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W. W. Pratt,[37] G. Pratten,[106] V. Predoi,[36] T. Prestegard,[20] M. Principe,[8] S. Privitera,[38] G. A. Prodi,[113,98] L. G. Prokhorov,[63] O. Puncken,[9,10] M. Punturo,[42] P. Puppo,[35] M. Pürrer,[38] H. Qi,[20] V. Quetschke,[108] E. A. Quintero,[1] R. Quitzow-James,[70] D. S. Rabeling,[24] H. Radkins,[46] P. Raffai,[54] S. Raja,[62] C. Rajan,[62] B. Rajbhandari,[84] M. Rakhmanov,[108] K. E. Ramirez,[108] A. Ramos-Buades,[106] Javed Rana,[18] P. Rapagnani,[100,35] V. Raymond,[36] M. Razzano,[21,22] J. Read,[28] T. Regimbau,[66,7] L. Rei,[61] S. Reid,[26] D. H. Reitze,[1,49] W. Ren,[11] F. Ricci,[100,35] P. M. Ricker,[11] K. Riles,[124] M. Rizzo,[59] N. A. Robertson,[1,45] R. Robie,[45] F. Robinet,[27] T. Robson,[105] A. Rocchi,[33] L. Rolland,[7] J. G. Rollins,[1] V. J. Roma,[70] R. Romano,[3,4] C. L. Romel,[46] J. H. Romie,[6] D. Rosińska,[157,55] M. P. Ross,[158] S. Rowan,[45] A. Rüdiger,[9,10] P. Ruggi,[30] G. Rutins,[159] K. Ryan,[46] S. Sachdev,[1] T. Sadecki,[46] M. Sakellariadou,[160] L. Salconi,[30] M. Saleem,[120] F. Salemi,[9] A. Samajdar,[150,13] L. Sammut,[5] L. M. Sampson,[91] E. J. Sanchez,[1] L. E. Sanchez,[1] N. Sanchis-Gual,[23] V. Sandberg,[46] J. R. Sanders,[43] N. Sarin,[5] B. Sassolas,[25] P. R. Saulson,[43] O. Sauter,[124] R. L. Savage,[46] A. Sawadsky,[34] P. Schale,[70] M. Scheel,[47] J. Scheuer,[91] P. Schmidt,[65] R. Schnabel,[34] R. M. S. Schofield,[70] A. Schönbeck,[34] E. Schreiber,[9,10] D. Schuette,[9,10] B. W. Schulte,[9,10] B. F. Schutz,[36,9] S. G. Schwalbe,[37] J. Scott,[45] S. M. Scott,[24] E. Seidel,[11] D. Sellers,[6] A. S. Sengupta,[161] D. Sentenac,[30] V. Sequino,[32,33,16] A. Sergeev,[130] Y. Setyawati,[9] D. A. Shaddock,[24] T. J. Shaffer,[46] A. A. Shah,[140] M. S. Shahriar,[91] M. B. Shaner,[112] L. Shao,[38] B. Shapiro,[50] P. Shawhan,[76] H. Shen,[11] D. H. Shoemaker,[14] D. M. Shoemaker,[77] K. Siellez,[77] X. Siemens,[20] M. Sieniawska,[55] D. Sigg,[46] A. D. Silva,[15] L. P. Singer,[80] A. Singh,[9,10] A. Singhal,[16,35] A. M. Sintes,[106] B. J. J. Slagmolen,[24] T. J. Slaven-Blair,[64] B. Smith,[6] J. R. Smith,[28] R. J. E. Smith,[5] S. Somala,[162] E. J. Son,[133] B. Sorazu,[45] F. Sorrentino,[61] T. Souradeep,[18] A. P. Spencer,[45] A. K. Srivastava,[110] K. Staats,[37] M. Steinke,[9,10] J. Steinlechner,[34,45] S. Steinlechner,[34] D. Steinmeyer,[9,10] B. Steltner,[9,10] S. P. Stevenson,[156] D. Stocks,[50] R. Stone,[108] D. J. Stops,[60] K. A. Strain,[45] G. Stratta,[72,73] S. E. Strigin,[63] A. Strunk,[46] R. Sturani,[163] A. L. Stuver,[164] T. Z. Summerscales,[165] L. Sun,[99] S. Sunil,[110] J. Suresh,[18] P. J. Sutton,[36] B. L. Swinkels,[13] M. J. Szczepańczyk,[37] M. Tacca,[13] S. C. Tait,[45] C. Talbot,[5] D. Talukder,[70] D. B. Tanner,[49] M. Tápai,[121] A. Taracchini,[38] J. D. Tasson,[95] J. A. Taylor,[140] R. Taylor,[1] S. V. Tewari,[154] T. Theeg,[9,10] F. Thies,[9,10] E. G. Thomas,[60] M. Thomas,[6] P. Thomas,[46] K. A. Thorne,[6] E. Thrane,[5] S. Tiwari,[16,98] V. Tiwari,[36] K. V. Tokmakov,[26] K. Toland,[45] M. Tonelli,[21,22] Z. Tornasi,[45] A. Torres-Forné,[23] C. I. Torrie,[1] D. Töyrä,[60] F. Travasso,[30,42] G. Traylor,[6] J. Trinastic,[49] M. C. Tringali,[113,98] L. Trozzo,[166,22] K. W. Tsang,[13] M. Tse,[14] R. Tso,[47] D. Tsuna,[82] L. Tsukada,[82] D. Tuyenbayev,[108] K. Ueno,[20] D. Ugolini,[167] A. L. Urban,[1] S. A. Usman,[36] H. Vahlbruch,[9,10] G. Vajente,[1] G. Valdes,[2] N. van Bakel,[13] M. van Beuzekom,[13] J. F. J. van den Brand,[75,13] C. Van Den Broeck,[13,168] D. C. Vander-Hyde,[43] L. van der Schaaf,[13] J. V. van Heijningen,[13] A. A. van Veggel,[45] M. Vardaro,[52,53] V. Varma,[47] S. Vass,[1] M. Vasúth,[48] A. Vecchio,[60] G. Vedovato,[53] J. Veitch,[45] P. J. Veitch,[56] K. Venkateswara,[158] G. Venugopalan,[1] D. Verkindt,[7] F. Vetrano,[72,73] A. Viceré,[72,73] A. D. Viets,[20] S. Vinciguerra,[60] D. J. Vine,[159] J.-Y. Vinet,[66] S. Vitale,[14] T. Vo,[43] H. Vocca,[41,42] C. Vorvick,[46] S. P. Vyatchanin,[63] A. R. Wade,[1] L. E. Wade,[85] M. Wade,[85] R. Walet,[13] M. Walker,[28] L. Wallace,[1] S. Walsh,[20,9] G. Wang,[16,22] H. Wang,[60] J. Z. Wang,[124] W. H. Wang,[108] Y. F. Wang,[96] R. L. Ward,[24] J. Warner,[46] M. Was,[7] J. Watchi,[101] B. Weaver,[46] L.-W. Wei,[9,10] M. Weinert,[9,10] A. J. Weinstein,[1] R. Weiss,[14] F. Wellmann,[9,10] L. Wen,[64] E. K. Wessel,[11] P. Weßels,[9,10] J. Westerweck,[9] K. Wette,[24] J. T. Whelan,[59] B. F. Whiting,[49] C. Whittle,[14] D. Wilken,[9,10] D. Williams,[45] R. D. Williams,[1] A. R. Williamson,[59,65] J. L. Willis,[1,125] B. Willke,[9,10] M. H. Wimmer,[9,10] W. Winkler,[9,10] C. C. Wipf,[1] H. Wittel,[9,10] G. Woan,[45] J. Woehler,[9,10] J. K. Wofford,[59] W. K. Wong,[96] J. Worden,[46] J. L. Wright,[45] D. S. Wu,[9,10] D. M. Wysocki,[59] S. Xiao,[1] W. Yam,[14] H. Yamamoto,[1] C. C. Yancey,[76] L. Yang,[170] M. J. Yap,[24] M. Yazback,[49] Hang Yu,[14] Haocun Yu,[14] M. Yvert,[7] A. Zadrożny,[135] M. Zanolin,[37] T. Zelenova,[30] J.-P. Zendri,[53] M. Zevin,[91] J. Zhang,[64] L. Zhang,[1] M. Zhang,[145] T. Zhang,[45] Y.-H. Zhang,[9,10] C. Zhao,[64] M. Zhou,[91] Z. Zhou,[91] S. J. Zhu,[9,10] X. J. Zhu,[5] A. B. Zimmerman,[92,171] M. E. Zucker,[1,14] and J. Zweizig[1]

(LIGO Scientific Collaboration and Virgo Collaboration)

N. N. Weinberg[169]

[1]LIGO, California Institute of Technology, Pasadena, California 91125, USA
[2]Louisiana State University, Baton Rouge, Louisiana 70803, USA
[3]Università di Salerno, Fisciano, I-84084 Salerno, Italy
[4]INFN, Sezione di Napoli, Complesso Universitario di Monte S.Angelo, I-80126 Napoli, Italy
[5]OzGrav, School of Physics & Astronomy, Monash University, Clayton 3800, Victoria, Australia
[6]LIGO Livingston Observatory, Livingston, Louisiana 70754, USA
[7]Laboratoire d'Annecy de Physique des Particules (LAPP), University Grenoble Alpes, Université Savoie Mont Blanc, CNRS/IN2P3, F-74941 Annecy, France
[8]University of Sannio at Benevento, I-82100 Benevento, Italy and INFN, Sezione di Napoli, I-80100 Napoli, Italy
[9]Max Planck Institute for Gravitational Physics (Albert Einstein Institute), D-30167 Hannover, Germany
[10]Leibniz Universität Hannover, D-30167 Hannover, Germany
[11]NCSA, University of Illinois at Urbana-Champaign, Urbana, Illinois 61801, USA
[12]University of Cambridge, Cambridge CB2 1TN, United Kingdom
[13]Nikhef, Science Park 105, 1098 XG Amsterdam, Netherlands
[14]LIGO, Massachusetts Institute of Technology, Cambridge, Massachusetts 02139, USA
[15]Instituto Nacional de Pesquisas Espaciais, 12227-010 São José dos Campos, São Paulo, Brazil
[16]Gran Sasso Science Institute (GSSI), I-67100 L'Aquila, Italy
[17]INFN, Laboratori Nazionali del Gran Sasso, I-67100 Assergi, Italy
[18]Inter-University Centre for Astronomy and Astrophysics, Pune 411007, India
[19]International Centre for Theoretical Sciences, Tata Institute of Fundamental Research, Bengaluru 560089, India
[20]University of Wisconsin-Milwaukee, Milwaukee, Wisconsin 53201, USA
[21]Università di Pisa, I-56127 Pisa, Italy
[22]INFN, Sezione di Pisa, I-56127 Pisa, Italy
[23]Departamento de Astronomía y Astrofísica, Universitat de València, E-46100 Burjassot, València, Spain
[24]OzGrav, Australian National University, Canberra, Australian Capital Territory 0200, Australia
[25]Laboratoire des Matériaux Avancés (LMA), CNRS/IN2P3, F-69622 Villeurbanne, France
[26]SUPA, University of Strathclyde, Glasgow G1 1XQ, United Kingdom
[27]LAL, University Paris-Sud, CNRS/IN2P3, Université Paris-Saclay, F-91898 Orsay, France
[28]California State University Fullerton, Fullerton, California 92831, USA
[29]APC, AstroParticule et Cosmologie, Université Paris Diderot, CNRS/IN2P3, CEA/Irfu, Observatoire de Paris, Sorbonne Paris Cité, F-75205 Paris Cedex 13, France
[30]European Gravitational Observatory (EGO), I-56021 Cascina, Pisa, Italy
[31]Chennai Mathematical Institute, Chennai 603103, India
[32]Università di Roma Tor Vergata, I-00133 Roma, Italy
[33]INFN, Sezione di Roma Tor Vergata, I-00133 Roma, Italy
[34]Universität Hamburg, D-22761 Hamburg, Germany
[35]INFN, Sezione di Roma, I-00185 Roma, Italy
[36]Cardiff University, Cardiff CF24 3AA, United Kingdom
[37]Embry-Riddle Aeronautical University, Prescott, Arizona 86301, USA
[38]Max Planck Institute for Gravitational Physics (Albert Einstein Institute), D-14476 Potsdam-Golm, Germany
[39]Korea Institute of Science and Technology Information, Daejeon 34141, Korea
[40]West Virginia University, Morgantown, West Virginia 26506, USA
[41]Università di Perugia, I-06123 Perugia, Italy
[42]INFN, Sezione di Perugia, I-06123 Perugia, Italy
[43]Syracuse University, Syracuse, New York 13244, USA
[44]University of Minnesota, Minneapolis, Minnesota 55455, USA
[45]SUPA, University of Glasgow, Glasgow G12 8QQ, United Kingdom
[46]LIGO Hanford Observatory, Richland, Washington 99352, USA
[47]Caltech CaRT, Pasadena, California 91125, USA
[48]Wigner RCP, RMKI, H-1121 Budapest, Konkoly Thege Miklós út 29-33, Hungary
[49]University of Florida, Gainesville, Florida 32611, USA
[50]Stanford University, Stanford, California 94305, USA
[51]Università di Camerino, Dipartimento di Fisica, I-62032 Camerino, Italy
[52]Università di Padova, Dipartimento di Fisica e Astronomia, I-35131 Padova, Italy
[53]INFN, Sezione di Padova, I-35131 Padova, Italy
[54]MTA-ELTE Astrophysics Research Group, Institute of Physics, Eötvös University, Budapest 1117, Hungary
[55]Nicolaus Copernicus Astronomical Center, Polish Academy of Sciences, 00-716, Warsaw, Poland
[56]OzGrav, University of Adelaide, Adelaide, South Australia 5005, Australia
[57]Dipartimento di Scienze Matematiche, Fisiche e Informatiche, Università di Parma, I-43124 Parma, Italy
[58]INFN, Sezione di Milano Bicocca, Gruppo Collegato di Parma, I-43124 Parma, Italy

[59]*Rochester Institute of Technology, Rochester, New York 14623, USA*
[60]*University of Birmingham, Birmingham B15 2TT, United Kingdom*
[61]*INFN, Sezione di Genova, I-16146 Genova, Italy*
[62]*RRCAT, Indore, Madhya Pradesh 452013, India*
[63]*Faculty of Physics, Lomonosov Moscow State University, Moscow 119991, Russia*
[64]*OzGrav, University of Western Australia, Crawley, Western Australia 6009, Australia*
[65]*Department of Astrophysics/IMAPP, Radboud University Nijmegen, P.O. Box 9010, 6500 GL Nijmegen, Netherlands*
[66]*Artemis, Université Côte d'Azur, Observatoire Côte d'Azur, CNRS, CS 34229, F-06304 Nice Cedex 4, France*
[67]*Physik-Institut, University of Zurich, Winterthurerstrasse 190, 8057 Zurich, Switzerland*
[68]*University of Rennes, CNRS, Institut FOTON—UMR6082, F-3500 Rennes, France*
[69]*Washington State University, Pullman, Washington 99164, USA*
[70]*University of Oregon, Eugene, Oregon 97403, USA*
[71]*Laboratoire Kastler Brossel, Sorbonne Université, CNRS, ENS-Université PSL, Collège de France, F-75005 Paris, France*
[72]*Università degli Studi di Urbino 'Carlo Bo,' I-61029 Urbino, Italy*
[73]*INFN, Sezione di Firenze, I-50019 Sesto Fiorentino, Firenze, Italy*
[74]*Astronomical Observatory Warsaw University, 00-478 Warsaw, Poland*
[75]*VU University Amsterdam, 1081 HV Amsterdam, Netherlands*
[76]*University of Maryland, College Park, Maryland 20742, USA*
[77]*School of Physics, Georgia Institute of Technology, Atlanta, Georgia 30332, USA*
[78]*Université Claude Bernard Lyon 1, F-69622 Villeurbanne, France*
[79]*Università di Napoli 'Federico II,' Complesso Universitario di Monte S.Angelo, I-80126 Napoli, Italy*
[80]*NASA Goddard Space Flight Center, Greenbelt, Maryland 20771, USA*
[81]*Dipartimento di Fisica, Università degli Studi di Genova, I-16146 Genova, Italy*
[82]*RESCEU, University of Tokyo, Tokyo, 113-0033, Japan*
[83]*Tsinghua University, Beijing 100084, China*
[84]*Texas Tech University, Lubbock, Texas 79409, USA*
[85]*Kenyon College, Gambier, Ohio 43022, USA*
[86]*The University of Mississippi, University, Mississippi 38677, USA*
[87]*Museo Storico della Fisica e Centro Studi e Ricerche "Enrico Fermi", I-00184 Roma, Italyrico Fermi, I-00184 Roma, Italy*
[88]*The Pennsylvania State University, University Park, Pennsylvania 16802, USA*
[89]*National Tsing Hua University, Hsinchu City, 30013 Taiwan, Republic of China*
[90]*Charles Sturt University, Wagga Wagga, New South Wales 2678, Australia*
[91]*Center for Interdisciplinary Exploration & Research in Astrophysics (CIERA), Northwestern University, Evanston, Illinois 60208, USA*
[92]*Canadian Institute for Theoretical Astrophysics, 60 St. George Street, Toronto, Ontario, M5S 3H8, Canada*
[93]*University of Chicago, Chicago, Illinois 60637, USA*
[94]*Pusan National University, Busan 46241, Korea*
[95]*Carleton College, Northfield, Minnesota 55057, USA*
[96]*The Chinese University of Hong Kong, Shatin, NT, Hong Kong*
[97]*INAF, Osservatorio Astronomico di Padova, I-35122 Padova, Italy*
[98]*INFN, Trento Institute for Fundamental Physics and Applications, I-38123 Povo, Trento, Italy*
[99]*OzGrav, University of Melbourne, Parkville, Victoria 3010, Australia*
[100]*Università di Roma 'La Sapienza,' I-00185 Roma, Italy*
[101]*Université Libre de Bruxelles, Brussels 1050, Belgium*
[102]*Sonoma State University, Rohnert Park, California 94928, USA*
[103]*Departamento de Matemáticas, Universitat de València, E-46100 Burjassot, València, Spain*
[104]*Columbia University, New York, New York 10027, USA*
[105]*Montana State University, Bozeman, Montana 59717, USA*
[106]*Universitat de les Illes Balears, IAC3—IEEC, E-07122 Palma de Mallorca, Spain*
[107]*University of Rhode Island, Kingston, RI 02881, USA*
[108]*The University of Texas Rio Grande Valley, Brownsville, Texas 78520, USA*
[109]*Bellevue College, Bellevue, Washington 98007, USA*
[110]*Institute for Plasma Research, Bhat, Gandhinagar 382428, India*
[111]*The University of Sheffield, Sheffield S10 2TN, United Kingdom*
[112]*California State University, Los Angeles, 5151 State University Dr, Los Angeles, California 90032, USA*
[113]*Università di Trento, Dipartimento di Fisica, I-38123 Povo, Trento, Italy*
[114]*Montclair State University, Montclair, New Jersey 07043, USA*
[115]*National Astronomical Observatory of Japan, 2-21-1 Osawa, Mitaka, Tokyo 181-8588, Japan*
[116]*Canadian Institute for Theoretical Astrophysics, University of Toronto, Toronto, Ontario M5S 3H8, Canada*
[117]*Observatori Astronòmic, Universitat de València, E-46980 Paterna, València, Spain*

[118]*School of Mathematics, University of Edinburgh, Edinburgh EH9 3FD, United Kingdom*
[119]*University and Institute of Advanced Research, Koba Institutional Area, Gandhinagar Gujarat 382007, India*
[120]*Indian Institute of Technology Bombay, Powai, Mumbai 400076 India*
[121]*University of Szeged, Dóm tér 9, Szeged 6720, Hungary*
[122]*Tata Institute of Fundamental Research, Mumbai 400005, India*
[123]*INAF, Osservatorio Astronomico di Capodimonte, I-80131, Napoli, Italy*
[124]*University of Michigan, Ann Arbor, Michigan 48109, USA*
[125]*Abilene Christian University, Abilene, Texas 79699, USA*
[126]*American University, Washington, DC, 20016, USA*
[127]*University of Białystok, 15-424 Białystok, Poland*
[128]*University of Southampton, Southampton SO17 1BJ, United Kingdom*
[129]*University of Washington Bothell, 18115 Campus Way NE, Bothell, Washington 98011, USA*
[130]*Institute of Applied Physics, Nizhny Novgorod, 603950, Russia*
[131]*Korea Astronomy and Space Science Institute, Daejeon 34055, Korea*
[132]*Inje University Gimhae, South Gyeongsang 50834, Korea*
[133]*National Institute for Mathematical Sciences, Daejeon 34047, Korea*
[134]*Ulsan Nationatl Institute of Science and Technology, Ulsan 44919, Korea*
[135]*NCBJ, 05-400 Świerk-Otwock, Poland*
[136]*Institute of Mathematics, Polish Academy of Sciences, 00656 Warsaw, Poland*
[137]*Cornell University, Ithaca, New York 14853, USA*
[138]*Hillsdale College, Hillsdale, Michigan 49242, USA*
[139]*Hanyang University, Seoul 04763, Korea*
[140]*NASA Marshall Space Flight Center, Huntsville, Alabama 35811, USA*
[141]*Dipartimento di Fisica, Università degli Studi Roma Tre, I-00154 Roma, Italy*
[142]*INFN, Sezione di Roma Tre, I-00154 Roma, Italy*
[143]*ESPCI, CNRS, F-75005 Paris, France*
[144]*Southern University and A&M College, Baton Rouge, Louisiana 70813, USA*
[145]*College of William and Mary, Williamsburg, Virginia 23187, USA*
[146]*Centre Scientifique de Monaco, 8 quai Antoine Ier, MC-98000, Monaco*
[147]*Indian Institute of Technology Madras, Chennai 600036, India*
[148]*INFN Sezione di Torino, Via P. Giuria 1, I-10125 Torino, Italy*
[149]*Institut des Hautes Etudes Scientifiques, F-91440 Bures-sur-Yvette, France*
[150]*IISER-Kolkata, Mohanpur, West Bengal 741252, India*
[151]*Whitman College, 345 Boyer Avenue, Walla Walla, Washington 99362 USA*
[152]*Scuola Normale Superiore, Piazza dei Cavalieri 7, I-56126 Pisa, Italy*
[153]*Université de Lyon, F-69361 Lyon, France*
[154]*Hobart and William Smith Colleges, Geneva, New York 14456, USA*
[155]*Università degli Studi di Firenze, I-50121 Firenze, Italy*
[156]*OzGrav, Swinburne University of Technology, Hawthorn VIC 3122, Australia*
[157]*Janusz Gil Institute of Astronomy, University of Zielona Góra, 65-265 Zielona Góra, Poland*
[158]*University of Washington, Seattle, Washington 98195, USA*
[159]*SUPA, University of the West of Scotland, Paisley PA1 2BE, United Kingdom*
[160]*King's College London, University of London, London WC2R 2LS, United Kingdom*
[161]*Indian Institute of Technology, Gandhinagar Ahmedabad Gujarat 382424, India*
[162]*Indian Institute of Technology Hyderabad, Sangareddy, Khandi, Telangana 502285, India*
[163]*International Institute of Physics, Universidade Federal do Rio Grande do Norte, Natal RN 59078-970, Brazil*
[164]*Villanova University, 800 Lancaster Ave, Villanova, Pennsylvania 19085, USA*
[165]*Andrews University, Berrien Springs, Michigan 49104, USA*
[166]*Università di Siena, I-53100 Siena, Italy*
[167]*Trinity University, San Antonio, Texas 78212, USA*
[168]*Van Swinderen Institute for Particle Physics and Gravity, University of Groningen, Nijenborgh 4, 9747 AG Groningen, Netherlands*
[169]*Department of Physics, and Kavli Institute for Astrophysics and Space Research, Massachusetts Institute of Technology, Cambridge, Massachusetts 02139, USA*
[170]*Colorado State University, Fort Collins, Colorado 80523, USA*
[171]*University of Texas at Austin, Austin, Texas 78712, USA*

[†]Deceased.